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SUSY-breaking Soft Terms in a MSSM Magnetized D7-brane Model

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Abstract

We compute the SUSY-breaking soft terms in a magnetized D7-brane model with MSSM-like spectrum, under the general assumption of non-vanishing auxiliary fields of the dilaton and Kähler moduli. As a particular scenario we discuss SUSY breaking triggered by ISD or IASD 3-form fluxes.

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1 Introduction

Recently a simple intersecting D-brane model was proposed with massless chiral spectrum close to that of the MSSM [1]. In this model the SM fields lie at the intersections of four sets of D6-branes wrapping an (orientifolded) toroidal compactification of Type IIA string theory. The same model may be equivalently described in terms of different T-dual configurations, e.g. in terms of a Type IIB orientifold with (magnetized) D9-branes and D5-branes [2]. Recently [3] it has been shown how this type of D-brane configurations may be promoted to a fully $\mathcal{N}$=1 SUSY tadpole-free model (see also [4]) by embedding it into a $\mathbb{Z}_2 \times \mathbb{Z}_2$ Type IIB orientifold along the lines suggested in [5]. A number of results for the effective Lagrangian in such type of D-brane models is known by now. The Yukawa couplings among chiral fields were computed in [1, 6, 7, 2] and other aspects of the effective action may be found in [8, 9, 10, 11, 12, 13, 14, 15]. For an up-to-date review, see [16].

One interesting point to address is the structure of possible SUSY-breaking soft terms in this model. It has been recently realized that fluxes of antisymmetric R-R and NS-NS fields in Type IIB orientifolds may provide a source of such terms [17, 18, 19, 20, 13, 21, 22]. It was also realized [18] that SUSY breaking from imaginary self-dual (ISD) 3-form fluxes correspond to a non-vanishing vev for the auxiliary field of the overall modulus T and imaginary anti-self-dual (IASD) correspond to a non vanishing vev for the auxiliary field of the complex dilaton S. On the other hand, a possible phenomenological application of these ideas was proposed in [21]. In particular, if one assumes that the SM particles correspond to geometric D7-brane moduli, a simple set of SUSY-breaking soft terms may be shown to arise from ISD fluxes.

In this article we would like to present explicit results for the SUSY breaking soft terms in the MSSM-like model of ref. [1] as a function of the vevs of the auxiliary fields of the Kähler moduli T_i and/or the complex dilaton S. As particular examples we consider vevs induced by ISD and/or IASD 3-form fluxes. We construct the model [1] in terms of 3 stacks of intersecting D7-branes, one of them containing a constant magnetic field (leading to chirality and family replication). We then use the effective supergravity Lagrangian

approach in order to obtain soft terms, as in ref. [23, 24].

A previous detailed analysis of these soft terms including the effect of the non-vanishing magnetic flux was presented by Lüst, Reffert and Stieberger in ref. [14], based on the Kähler metrics of matter fields computed in ref. [11, 12]. These included the effect of magnetic fluxes. Some phenomenological analysis of those results was described in [25]. Soft terms in the T-dominance case were briefly discussed in [21] following [26], in which the effect of magnetic fluxes was not included. In the present paper we revisit previous results taking into account a proper normalization of the matter fields. We also make use of the Kähler metrics for chiral fields discussed in ref. [11, 12]. Including a factor implied by the analysis of [11] the SUSY-breaking soft terms simplify considerably. One of the motivations of the present work was to find the connection with the analogous results obtained in [26] in the absence of magnetic fluxes. Indeed, we find that in the limit of diluted magnetic fluxes those results are recovered.

It is known that NS-NS and R-R fluxes on *toroidal* settings induce soft terms on D3-brane fields of order M_s^2/M_{Pl}, so that one can obtain a hierarchy of scales by lowering the string scale [18]. However, in the case of intersecting D7-branes, as in the model at hand, one cannot lower the string scale without making the SM gauge couplings unacceptably small. Therefore, in toroidal/orbifold models with intersecting D7-branes the fluxed-induced soft terms are typically of order the string scale. This fact is due to the simplicity of toroidal compactifications in which the compact space is flat and the fluxes are distributed uniformly. In a generic Calabi-Yau (CY) compactification this is not going to be the case and there may be regions in the CY in which fluxes are concentrated and others in which fluxes are diluted. This possibility was considered e.g. in [27, 28] in order to obtain hierarchies. Thus, for generic CY compactifications the size of soft terms will actually depend on the detailed geometry of the fluxes in the CY.

The local set of branes leading to a MSSM-like spectrum introduced in [1] is nevertheless likely to be more generic than the toroidal setting in which it was first proposed (see e.g. [29]). In particular, it has recently been shown [30] that there are many thousands of models with the 4-stacks of branes structure of the model in [1] (these are labeled Type-4 models in ref. [30]). Therefore, one may expect to obtain this MSSM structure in CY

orientifold models beyond the toroidal setting. In these more general models the size of induced soft terms may not be tied to the string scale and could be much lower. In our effective field theory analysis below we will not commit ourselves to a particular scale for the soft terms. Instead, following [23, 24], we will assume that the effect of SUSY-breaking is encoded in non-vanishing vevs for the auxiliary fields of the complex dilaton and Kähler moduli. However, we also discuss the case in which the source of SUSY-breaking are constant IASD or/and IASD 3-form fluxes, which correspond to a particular choice for the auxiliary fields.

2 A MSSM-like model from magnetized D7-branes

We will construct the model in [1] in terms of three sets of intersecting Type IIB D7-branes (see e.g.[3]). We consider type IIB string theory compactified on a factorized six-torus $\mathrm{T}^6 = \otimes_{i=1}^3 \mathrm{T}_i^2$. We will further do an orientifold projection by $\Omega(-1)^F I_6$, Ω being the world-sheet parity operator and I_6 a simultaneous reflexion of the six toroidal coordinates. We also include sets of $\mathrm{D}9_a$-branes and allow for possible constant magnetic fluxes across any of the three 2-tori

$$\frac{m_a^i}{2\pi} \int_{\mathrm{T}_i^2} F_a^i = n_a^i \ , \tag{2.1}$$

where F_a^i is the world-volume magnetic field. For each group of branes the state of magnetization is thus characterized by the integers (n_a^i, m_a^i), where m_a^i is the wrapping number and n_a^i is the total magnetic flux. It is useful to introduce the angles

$$\psi_a^i = \arctan 2\pi\alpha' F_a^i = \arctan \frac{\alpha' n_a^i}{m_a^i A_i} \tag{2.2}$$

where $(2\pi)^2 A_i$ is the area of the T_i^2. The magnetized $\mathrm{D}9_a$-branes [31, 32, 33] preserve the same supersymmetry of the orientifold planes provided that [34, 35]

$$\sum_{i=1}^{3} \psi_a^i = \frac{3\pi}{2} \bmod 2\pi \ . \tag{2.3}$$

Note that in this scheme lower dimensional branes are described setting $m_a^i = 0$ for all i transverse to the brane. For example, a D3-brane has $(n_a^i, m_a^i) = (1, 0)$, $i = 1, 2, 3$. Notice

in particular that the D3-brane satisfies (2.3). T-duality along the horizontal direction in each T_i^2 gives the dual picture of D6-branes at angles. For example, for square T_i^2, $A_i = R_{ix}R_{iy}$, and the dual angle is $\vartheta_a^i = \arctan(n_a^i R_{ix}/m_a^i R_{iy})$.

In order to reproduce the structure of the MSSM-like model of ref. [1, 2] one introduces three sets of D7$_i$-branes $i = 1, 2, 3$ which are characterized by being transverse to the i-th 2-torus. In particular the relevant magnetic data is

$$\begin{array}{ccccc} \text{Branes} & (n_a^1, m_a^1) & (n_a^2, m_a^2) & (n_a^3, m_a^3) & (\psi_a^1, \psi_a^2, \psi_a^3) \\ \mathrm{D7}_1 & (1{,}0) & (g, 1) & (g, -1) & (\frac{\pi}{2}, \pi\delta_2, \pi - \pi\delta_3) \\ \mathrm{D7}_2 & (0{,}1) & (1, 0) & (0, -1) & (0, \frac{\pi}{2}, \pi) \\ \mathrm{D7}_3 & (0{,}1) & (0, -1) & (1, 0) & (0, \pi, \frac{\pi}{2}) \end{array} . \tag{2.4}$$

where $\pi\delta_i = \arctan(\alpha' g/A_i)$. We will take D7$_1$-branes to come in four copies so that generically the associated gauge group will be $U(4)$. Branes D7$_2$ and D7$_3$ come only in one copy and are located on top of the orientifold plane at the origin so that they give rise to a gauge group $Sp(2) \times Sp(2) \simeq SU(2) \times SU(2)$. Altogether the overall gauge group is $U(4) \times SU(2) \times SU(2)$. It may be shown that the $U(1)$ (which corresponds to $(3B + L)$) is anomalous and becomes massive in the usual way by the Stückelberg mechanism. In fact, one can further make the breakings $SU(4) \to SU(3)_c \times U(1)_{B-L}$ and $SU(2)_R \to U(1)_R$ by e.g. Wilson lines. Thus the final gauge group is just $SU(3) \times SU(2) \times U(1)_Y \times U(1)_{B-L}$. The final chiral spectrum is displayed in Table 1 for the choice $g = 3$ which leads to three quark/lepton generations.

The D7$_2$ and D7$_3$ do not have magnetic flux and do verify the supersymmetric condition (2.3). However, for the case of the D7$_1$-branes, with opposite magnetic fields turned on in the second and third T^2, to be supersymmetric we need to impose

$$A_2 = A_3 = A \ , \tag{2.5}$$

so that $\delta_2 = \delta_3 = \delta$ and

$$\tan \pi\delta = \frac{\alpha' g}{A}. \tag{2.6}$$

Clearly, the condition (2.3) also guarantees that any two sets of branes preserve a common supersymmetry. Notice that the relative angles

$$\theta_{ab}^i = \psi_b^i - \psi_a^i \tag{2.7}$$

Intersection	Matter fields	Rep.	Q_{B-L}	Y
$\mathrm{D7}_1 - \mathrm{D7}_2$	Q_L	$3(3,2)$	1	1/6
$\mathrm{D7}_1 - \mathrm{D7}_3$	U_R	$3(\bar{3},1)$	-1	-2/3
$\mathrm{D7}_1 - \mathrm{D7}_3$	D_R	$3(\bar{3},1)$	-1	1/3
$\mathrm{D7}_1 - \mathrm{D7}_2$	E_L	$3(1,2)$	-1	1/2
$\mathrm{D7}_1 - \mathrm{D7}_3$	E_R	$3(1,1)$	1	-1
$\mathrm{D7}_1 - \mathrm{D7}_3$	N_R	$3(1,1)$	1	0
$\mathrm{D7}_2 - \mathrm{D7}_3$	H	$(1,2)$	0	1/2
$\mathrm{D7}_2 - \mathrm{D7}_3$	$\bar{H}$	$(1,2)$	0	-1/2

Table 1: Chiral spectrum of the MSSM-like model.

automatically satisfy $\sum_i \theta^i_{ab} = 0 \bmod 2\pi$.

Departures from the equality $A_2 = A_3$ may be shown [6, 8, 36, 37] to correspond to a non-vanishing Fayet-Iliopoulos term for the anomalous $U(1)_{3B+L}$. In fact, for $A_2 = (A_3 + \epsilon)$ and small ϵ one finds [8, 36]

$$\xi_{FI} \;=\; \frac{g\epsilon}{A^2 + \alpha'^2 g^2} \tag{2.8}$$

where the D-term potential is of the form

$$V_{FI}(\phi_n) = \frac{1}{2g^2_{U(1)}}\Big(\sum_n q_n|\phi_n|^2 + \xi_{FI}\Big)^2. \tag{2.9}$$

and ϕ_n runs over squarks and sleptons. Left-handed and right-handed chiral fields have positive and negative $U(1)_{3B+L}$ charge respectively so that a non-vanishing ϵ may induce further symmetry breaking. Note that this potential as it stands does not prefer $\xi_a = 0$ (and hence the SUSY condition $A_2 = A_3$) as sometimes claimed in the literature, since a non-vanishing ξ_a may always be compensated with a vev for a right-handed scalar field (e.g. the right-handed sneutrino). On the other hand, in the presence of soft masses for the chiral fields, as shown to appear in the next sections, a vanishing FI-term is dynamically preferred and so is the SUSY condition $A_2 = A_3$. As we will see, this implies in turn the unification of $SU(2)_L$ and $SU(2)_R$ gauge couplings.

Let us finally comment that, as it stands, this brane configuration has R-R tadpoles so some additional ('hidden') brane system should be added. This can be done in a way consistent with $\mathcal{N}=1$ SUSY if we embed this brane configuration in a $\mathbb{Z}_2\times\mathbb{Z}_2$ orientifold [5] as recently shown in [3] (although in this case one cannot do the breaking $SU(2)_R \to U(1)_R$ via Wilson lines [3]). Since we are only interested in the structure of soft terms for the MSSM fields we will not deal here with these global issues of the compactification. Our results will still hold for those global generalizations.

3 Massless fields and effective supergravity action

Let us now turn to the effective supergravity action in this model. We will compile general formulas for the Kähler potential, matter metrics and gauge kinetic function and we will apply them to the specific D-brane model at hand.

We begin with the field content. In the closed string sector, in addition to the supergravity multiplet, one has the dilaton S plus the Kähler and complex structure moduli. We use conventions such that the complex dilaton is given by

$$S = e^{-\phi_{10}} + i a_0 \ , \tag{3.1}$$

where a_0 is the R-R 0-form. Recall that the string coupling constant is $g_s = e^{\phi_{10}}$.

For the metric moduli we will restrict for simplicity here to the diagonal fields U_j, T_j, $j = 1, 2, 3$. Note that in any case the off-diagonal fields would not be present in a $\mathbb{Z}_2 \times \mathbb{Z}_2$ embedding of the present model. To be more concrete, let the T^2_j lattice vectors be denoted e_{jx}, e_{jy}. Then, the geometric toroidal moduli are

$$\begin{aligned} \tau_j &= \frac{1}{e_{jx}^2}(A_j + i\, e_{jx}\cdot e_{jy}) \\ \rho_j &= A_j + i a_j \ , \end{aligned} \tag{3.2}$$

where the axions a_j arise from the R-R 4-form. In type IIB, the toroidal complex structure τ_j is equal to the moduli field U_j that appears in the D=4 supergravity action. However, the correct Kähler moduli field T_j is not ρ_j. One way to see this is to realize that the gauge coupling squared of *unmagnetized* D7$_j$-branes should be equal to $2\pi/\mathrm{Re}\, T_j$.

Then, since, e.g. a D7_1 wraps T_2^2 and T_3^2, from the Born-Infeld action it follows that $\mathrm{Re}\, T_1 = e^{-\phi_{10}} A_2 A_3/\alpha'^2$. In general

$$T_i = e^{-\phi_{10}} \frac{A_j A_k}{\alpha'^2} + i a_i \quad ; \quad j \neq k \neq i \ . \tag{3.3}$$

For later convenience we define

$$s = S + \overline{S} \quad ; \quad t_i = T_i + \overline{T}_i \quad ; \quad u_i = U_i + \overline{U}_i \ . \tag{3.4}$$

The D=4 gravitational coupling is $G_N = \kappa^2/8\pi$ where

$$\kappa^{-2} = \frac{M_{Pl}^2}{8\pi} = e^{-2\phi_{10}} \frac{A_1 A_2 A_3}{\pi \alpha'^4} = \frac{(s t_1 t_2 t_3)^{1/2}}{4\pi\alpha'} \ . \tag{3.5}$$

It is also useful to introduce the T-duality invariant four-dimensional dilaton, namely

$$\phi_4 = \phi_{10} - \frac{1}{2} \log(A_1 A_2 A_3/\alpha'^3) \ . \tag{3.6}$$

Notice that $\kappa^{-2} = e^{-2\phi_4}/\pi\alpha'$. The string scale is $M_s = 1/\sqrt{\alpha'}$.

Open strings give rise to charged fields. We call 'untwisted' the states corresponding to open strings beginning and ending on the same stack of branes, whereas 'twisted' refers to the chiral fields lying at the intersection of two different stacks of D7-branes. For the content of branes in (2.4), and assuming supersymmetry is preserved, the untwisted sectors D7_i-D7_i, $i = 1, 2, 3$, give a gauge multiplet of a group G_i and 3 massless chiral multiplets, denoted $C_j^{7_i}$, $j = 1, 2, 3$, transforming in the adjoint of G_i. The $C_j^{7_i}$ are the D7_i-brane moduli, $C_i^{7_i}$ gives the position of the brane in the transverse T_i^2 whereas $C_j^{7_i}$, $j \neq i$, correspond to Wilson lines on the two internal complex dimensions parallel to the D7_i-brane. From the twisted sectors D7_i-D7_j there are only chiral massless multiplets, denoted $C^{7_i 7_j}$, transforming as bifundamentals of $G_i \times G_j$.

The low-energy dynamics of the massless fields is governed by a D=4, $\mathcal{N}$=1 supergravity action that depends on the Kähler potential, the gauge kinetic functions and the superpotential. In particular, the F-part of the scalar potential is

$$V = e^{\kappa^2 K} \left[K^{\bar{A}B} (D_A W)^* (D_B W) - 3\kappa^2 |W|^2 \right] \ , \tag{3.7}$$

where $D_A W = \partial_A W + \kappa^2 \partial_A K W$ and $K^{\bar{A}B}$ is the inverse of $K_{\bar{A}B} = \partial_{\bar{A}} \partial_B K$. Recall that the auxiliary field of a chiral superfield Φ_A is

$$\bar{F}^{\bar{A}} = \kappa^2 e^{\kappa^2 K/2} K^{\bar{A}B} D_B W \ . \tag{3.8}$$

We now describe the functions $K(\phi;\bar{\phi})$, $f_i(\phi)$ and $W(\phi)$ in our setup, in which $\Phi_A = \{M, C_I\}$, with $M = \{S, T_i, U_i\}$ and $C_I = \{C_j^{7_i}, C^{7_i 7_j}\}$.

3.1 Kähler potential

The Kähler potential has the structure

$$K = \hat{K}(M,\bar{M}) + \sum_{I,J} \tilde{K}_{I\bar{J}}(M,\bar{M}) C_I \bar{C}_J + \frac{1}{2}\sum_{I,J}[Z_{IJ}(M,\bar{M})C_I C_J + c.c.] + \cdots \ . \tag{3.9}$$

The contribution of the closed string moduli is

$$\kappa^2 \hat{K}(M,\bar{M}) = -\log s - \sum_{i=1}^{3}\log t_i - \sum_{i=1}^{3}\log u_i \ . \tag{3.10}$$

Below we describe in detail the Kähler metrics of matter fields.

The $\tilde{K}_{I\bar{J}}$ for unmagnetized branes were deduced in [26] using T-duality arguments. For generic magnetized branes they have been obtained in [11, 12] from a computation of string scattering amplitudes. These metrics vanish when $J \neq I$. Below we present the diagonal entries for all possible cases with the brane content of (2.4). To streamline notation we write $\tilde{K}_{i,j\bar{j}} = \tilde{K}_{C_j^{7_i}\bar{C}_j^{7_i}}$ in untwisted sectors, and $\tilde{K}_{ij,C\bar{C}} = \tilde{K}_{C^{7_i7_j}\bar{C}^{7_i7_j}}$ in twisted sectors. For untwisted fields we have

- D7_3-D7_3 (untwisted, unmagnetized)

$$\kappa^2\tilde{K}_{3,1\bar{1}} = \frac{1}{u_1 t_2} \quad ; \quad \kappa^2\tilde{K}_{3,2\bar{2}} = \frac{1}{u_2 t_1} \quad ; \quad \kappa^2\tilde{K}_{3,3\bar{3}} = \frac{1}{u_3 s} \ . \tag{3.11}$$

- D7_2-D7_2 (untwisted, unmagnetized)

$$\kappa^2\tilde{K}_{2,1\bar{1}} = \frac{1}{u_1 t_3} \quad ; \quad \kappa^2\tilde{K}_{2,2\bar{2}} = \frac{1}{u_2 s} \quad ; \quad \kappa^2\tilde{K}_{2,3\bar{3}} = \frac{1}{u_3 t_1} \ . \tag{3.12}$$

- D7_1-D7_1 (untwisted, magnetized)

$$\kappa^2\tilde{K}_{1,1\bar{1}} = \frac{1}{u_1 t_1 s}(g^2 s + t_1) \quad ; \quad \kappa^2\tilde{K}_{1,2\bar{2}} = \frac{1}{u_2 t_2} \quad ; \quad \kappa^2\tilde{K}_{1,3\bar{3}} = \frac{1}{u_3 t_3} \ . \tag{3.13}$$

To obtain these results we start from the general expressions in the geometric basis given in [12]. In our notation these are

$$\kappa^2\tilde{K}_{1,1\bar{1}} \;=\; \frac{e^{\phi_4}}{\alpha'^2\, u_1}\sqrt{\frac{\alpha' A_1}{A_2 A_3}}|A_3 m_1^3 + i\alpha' n_1^3||A_2 m_1^2 + i\alpha' n_1^2| \ ,$$

$$
\begin{aligned}
\kappa^2 \tilde{K}_{1,2\bar{2}} &= \frac{e^{\phi_4}}{u_2}\sqrt{\frac{\alpha' A_2}{A_1 A_3}}\left|\frac{A_3 m_1^3 + i\alpha' n_1^3}{A_2 m_1^2 + i\alpha' n_1^2}\right| \ , \\
\kappa^2 \tilde{K}_{1,3\bar{3}} &= \frac{e^{\phi_4}}{u_3}\sqrt{\frac{\alpha' A_3}{A_1 A_2}}\left|\frac{A_2 m_1^2 + i\alpha' n_1^2}{A_3 m_1^3 + i\alpha' n_1^3}\right| \ .
\end{aligned}
\tag{3.14}
$$

We then substitute the values of the (n_1^i, m_1^i) given in (2.4), use the supersymmetry condition (2.5) and also

$$
\alpha'^2 t_1 = sA^2 \tag{3.15}
$$

that follows from (3.3). When $m_1^i = 0$, (3.14) gives the metric of a D3-D3 sector. When $n_1^i = 0$, we just obtain the metric of unmagnetized D7$_1$-D7$_1$.

For the metrics of twisted fields one has

- D7$_2$-D7$_3$ (twisted, unmagnetized)

$$
\kappa^2 \tilde{K}_{23,C\bar{C}} = \frac{1}{(u_2 u_3 s t_1)^{1/2}} \ . \tag{3.16}
$$

- D7$_1$-D7$_2$ (twisted, magnetized)

$$
\kappa^2 \tilde{K}_{12,C\bar{C}} = \frac{1}{(stu_1)^{1/2} u_2^{1/2+\delta} u_3^{1-\delta}} \frac{\Gamma(\frac{1}{2}-\delta)}{\Gamma(1-\delta)} \ , \tag{3.17}
$$

where $t = t_2 = t_3$. Notice that δ depends implicitly on s and t_1. From (2.6) and (3.15),

$$
\tan \pi\delta = g(s/t_1)^{1/2} \ . \tag{3.18}
$$

Observe that $0 \leq \delta < \frac{1}{2}$. To derive (3.17) we start from

$$
\kappa^2 \tilde{K}_{12,C\bar{C}} = e^{\phi_4} \prod_{j=1}^{3} u_j^{-\nu_j} \sqrt{\frac{\Gamma(1-\nu_j)}{\Gamma(\nu_j)}} \ , \tag{3.19}
$$

where the ν_j, computed from $\hat{\nu}_j = \theta_{12}^j/\pi$, are such that $0 \leq \nu_j < 1$ and $\nu_1 + \nu_2 + \nu_3 = 2$. To determine the ν_j, observe first that $\hat{\nu}_1 + \hat{\nu}_2 + \hat{\nu}_3 = 0$. Assuming $\hat{\nu}_j \neq 0$ then implies that one or two of the $\hat{\nu}_j$ are negative. In the first case start instead from $\hat{\nu}_j = \theta_{21}^j/\pi$. Now two of the $\hat{\nu}_j$ are negative by construction. Finally, define $\nu_j = 1 + \hat{\nu}_j$ if $\hat{\nu}_j$ is negative, otherwise $\nu_j = \hat{\nu}_j$. In this case $\nu = (\frac{1}{2}, \frac{1}{2}+\delta, 1-\delta)$.

To arrive at (3.17) we use (3.18) and the relation

$$
\frac{\Gamma(\delta)}{\Gamma(\frac{1}{2}+\delta)} = \frac{\Gamma(\frac{1}{2}-\delta)}{\tan \pi\delta \, \Gamma(1-\delta)} \ . \tag{3.20}
$$

In (3.17) we can take the limit $\delta \to 0$ and recover the metric of unmagnetized D7_1-D7_2, provided we drop u_3 that would have exponent -1. Using again (3.18) and (3.20) we can also take the limit $\delta \to \frac{1}{2}$ and, dropping u_2 now with exponent -1, retrieve the metric of D3-D7_2, namely $\kappa^2 \tilde{K}_{C\bar{C}} = (t_1 t_3 u_1 u_3)^{-1/2}$. The fact that moduli u_j with would be exponent -1 do not appear in the metric also occurs in twisted sectors of heterotic orbifolds [40].

Eq. (3.19), including the prefactor e^{ϕ_4}, follows putting together results found in [11]. In the field basis this prefactor can be recast as $2(st_1t_2t_3)^{-1/4}$. The square root of Gamma functions, with the arguments as shown in (3.19), is determined by the differential equation that dictates the dependence on the Kähler moduli [11]. Finally, for the remaining twisted sector the metric is

- D7_1-D7_3 (twisted, magnetized)

$$\kappa^2 \tilde{K}_{13,C\bar{C}} = \frac{1}{(stu_1)^{1/2} u_2^{1-\delta}\, u_3^{1/2+\delta}} \frac{\Gamma(\frac{1}{2}-\delta)}{\Gamma(1-\delta)} \ . \tag{3.21}$$

3.2 Gauge kinetic functions

The gauge kinetic functions f_i for the groups arising in the D7_i-D7_i sectors are

$$f_1 = T_1 + g^2 S \quad ; \quad f_2 = T_2 \quad ; \quad f_3 = T_3 \ . \tag{3.22}$$

In general [8, 11],

$$\mathrm{Re}\, f_i = \frac{e^{-\phi_{10}}}{\alpha'^2} \prod_{j \neq i} |m_i^j A_j + i\alpha' n_i^j| \ . \tag{3.23}$$

Substituting the values of the (n_i^j, m_i^j) given in (2.4) leads to (3.22).

Note that if the SUSY condition (2.5) is verified, one has $\mathrm{Re}\, f_2 = \mathrm{Re}\, f_3$ and the $SU(2)_L$ and $SU(2)_R$ gauge couplings are unified. Note also that if the complete model has additional branes (as in e.g., [3]) the SUSY conditions may involve in general also the area A_1 of the first torus and imply further unification constraints.

Concerning the axions, one can check that the linear combination $(a_2 - a_3)$ becomes massive combining with the anomalous $U(1)_{3B+L}$ gauge boson through the Stückelberg mechanism. On the other hand, one can also check that the linear combination $(9a_0 - a_1)$ has axionic couplings with the QCD gauge bosons. This may help in solving the strong CP problem.

Further inspection of the f_i reveals an interesting bound on the string coupling constant in the D-brane model. From (3.18) and the $SU(4)$ gauge coupling α_1 we deduce the relation

$$\sin^2 \pi\delta = \frac{2\alpha_1 g^2}{g_s} \ . \tag{3.24}$$

This has a number of consequences. To have three generations, $g = 3$. Hence, the above relation implies $\alpha_1 \leq g_s/18$. For $\alpha_1(M_s) \sim 1/24$ this is consistent with $g_s < 1$. However, we already see that to get values of coupling constants of the order of (extrapolated) known gauge couplings, the string coupling constant g_s approaches the non-perturbative regime. In fact, in the present specific toroidal model in addition to the chiral spectrum there are massless chiral adjoints which will make the gauge interactions asymptotically non-free. Thus, the α_i will be larger than $\sim 1/24$ and hence the above bound will give $g_s > O(1)$. A second (related) implication of eq.(3.24) is that in order to accommodate gauge couplings consistent with experiment but still stay within the string perturbative regime with $g_s < 1$, the value of δ (and hence the magnetic flux) will be substantial. Thus, e.g. for $\alpha_1 = 1/24$ one finds $\pi\delta \sim 60^0$. All this tells us that in order to have consistency with the observed values of gauge couplings in this class of intersecting D7-brane models we would probably need to go to a non-perturbative F-theory (or perhaps simply non-toroidal) versions of them. Constructing such class of F-theory models would be a rather non-trivial task. In what follows we will not further deal with these issues and simply assume that we still remain in a perturbative regime, hoping that a proper fit of experimentally measured gauge couplings does not substantially modify our soft term results.

3.3 Superpotential

The superpotential can be written as

$$W(M, C_I) = \hat{W}(M) + \frac{1}{2}\sum_{I,J} \mu_{IJ}(M) C_I C_J + \frac{1}{6}\sum_{I,J,L} Y_{IJL}(M) C_I C_J C_L + \cdots \ . \tag{3.25}$$

For the brane content of (2.4), the cubic couplings allowed are of the form [38]

$$C_1^{7_i} C_2^{7_i} C_3^{7_i} \quad ; \quad d_{ijk} C_i^{7_j} C^{7_j 7_k} C^{7_j 7_k} \quad ; \quad C^{7_1 7_2} C^{7_2 7_3} C^{7_3 7_1} \ , \tag{3.26}$$

where $d_{ijk} = 1$ if $i \neq j \neq k$, otherwise $d_{ijk} = 0$. Note that for the case at hand the first type of couplings in (3.26) correspond to standard $\mathcal{N}$=4 Yukawa couplings among adjoints. Concerning the second type of couplings, only one of type $XH\bar{H}$ is present in the model. Here X is a linear combination of C_1^2, C_1^3, the latter corresponding to Wilson line chiral fields of the branes $D7_2$, $D7_3$ in the first complex plane. Note that a vev for X would render the Higgs multiplets massive so X behaves as a μ-term in the effective Lagrangian[1]. Finally, the third type of superpotential couplings corresponds to the regular Yukawa couplings between the chiral generations and the Higgs multiplets. All in all, the perturbative superpotential among the chiral open string multiplets in this model has the general expression (we write it for the extended gauge group $SU(4) \times SU(2)_L \times SU(2)_R$ for simplicity of notation)

$$W_{Yukawa} \; = \; \sum_i C_1^{7_i} C_2^{7_i} C_3^{7_i} \; + \; X\mathcal{H}\mathcal{H} \; + \; \sum_{\alpha,\beta} h_{\alpha\beta}\mathcal{H} L_\alpha R_\beta \; , \tag{3.27}$$

where $\mathcal{H}$, L_α and R_β are the Higgs and chiral fermions transforming as $(1,2,2)$, $(4,2,1)$ and $(\bar{4},1,2)$ respectively ($\alpha, \beta = 1,2,3$ are generation labels).

The superpotential couplings $h_{\alpha\beta}$ for this toroidal model have been computed in [1, 2] and are given by

$$h_{\alpha\beta} \; = \; \vartheta \begin{bmatrix} \frac{\alpha}{3} \\ 0 \end{bmatrix} (3\zeta_2, 3U_2)\, \vartheta \begin{bmatrix} \frac{\beta}{3} \\ 0 \end{bmatrix} (3\zeta_3, 3U_3) \;\; , \tag{3.28}$$

where ζ_i are certain combinations of singlet $C_i^{7_j}$ fields (see [1, 2]) and ϑ are Jacobi theta functions. For our purposes the only relevant thing to point out is that these superpotential couplings only depend on the complex structure moduli U_2, U_3, and not on the Kähler moduli nor the dilaton.

In principle we can use the above results to compute soft terms under the general assumption that the auxiliary fields of the moduli and dilaton are non-vanishing, in the spirit of refs.[23, 24]. We would not need then to specify the microscopic source of this values. On the other hand, lately we have learned that such microscopic source of SUSY-breaking may be provided by fluxes in Type II theory. Hence in addition to the above perturbative chiral couplings, a moduli dependent superpotential $\hat{W}(M)$ may be present.

[1] In the T-dual version in terms of intersecting D6-branes, $\langle X \rangle$ corresponds to the distance between the $SU(2)_L$ and $SU(2)_R$ D6-branes in the first complex plane.

In particular, it is known that antisymmetric R-R and NS-NS fluxes F_3, H_3 generate a superpotential [39]

$$W_f = \int G_3 \wedge \Omega \tag{3.29}$$

that depends on S and the U_i. Here $G_3 = F_3 - iSH_3$ and Ω is the holomorphic (3,0) form of T^6. Besides, there may be non-perturbative interactions (like e.g. those from gaugino condensation) giving rise to a generic superpotential W_{np}. Thus the total moduli-dependent superpotential will have the general form

$$\hat{W}(M) = W_f(S, U_i) + W_{np}(S, U_i, T_i) \ . \tag{3.30}$$

In the absence of W_{np} the equations of motion require G_3 to be imaginary self-dual (ISD), meaning that G_3 is a combination of (0,3) and (2,1) fluxes [27]. In this case $D_S W_f = 0$ and $D_{U_i} W_f = 0$ but $D_{T_i} W_f \neq 0$ because W_f does not depend on the T_i and a (0,3) piece in G_3 generates $W_f \neq 0$.

4 Soft Terms

Armed with all the above data for the low-energy effective action we can now compute the SUSY-breaking soft terms. To this purpose we will follow the approach in [23, 24] and assume that the auxiliary fields F_{T_i}, F_S of the Kähler moduli and the complex dilaton acquire non-vanishing expectation values. We will later consider the particular case in which the microscopic origin of such non-vanishing values is provided by ISD and IASD three-form fluxes. The standard results for the normalized soft parameters may be found e.g. in [24] and read

$$\begin{aligned} M_i &= \frac{1}{2\mathrm{Re}\, f_i} F^M \partial_M f_i \ , \\ m_I^2 &= m_{3/2}^2 + V_0 - \sum_{M,N} \bar{F}^{\bar{M}} F^N \partial_{\bar{M}} \partial_N \log(\tilde{K}_{I\bar{I}}) \ , \\ A_{IJL} &= F^M [\hat{K}_M + \partial_M \log(Y_{IJL}) - \partial_M \log(\tilde{K}_{I\bar{I}} \tilde{K}_{J\bar{J}} \tilde{K}_{L\bar{L}})] \ . \end{aligned} \tag{4.1}$$

Here V_0 is the vev of the scalar potential and the gravitino mass is

$$m_{3/2} = e^{\kappa^2 K/2} |W| \ . \tag{4.2}$$

Note that, as pointed out above, the Y_{IJL} superpotential couplings do not depend on the S and T_i fields, so that the second contribution to A_{IJL} in (4.1) vanishes identically. The expressions (4.1) are valid when $\tilde{K}_{I\bar{J}} \propto \delta_{I\bar{J}}$ which is our case.

The vevs of the auxiliary fields are conveniently parametrized as [24]

$$\begin{aligned} F^S &= \sqrt{3} s C m_{3/2} \sin\theta e^{-i\gamma_S} \\ F^{T_i} &= \sqrt{3} t_i \eta_i C m_{3/2} \cos\theta e^{-i\gamma_i} \ , \end{aligned} \tag{4.3}$$

where the goldstino angle θ and the η_i, with $\sum_i \eta_i^2 = 1$, control whether S or the T_i dominate SUSY breaking. We further assume that $F^{U_i} = 0$. Then, substituting in (3.7) gives

$$C^2 = 1 + \frac{V_0}{3m_{3/2}^2} \ . \tag{4.4}$$

We will now present the soft terms for the D-brane model.

<u>Gaugino masses</u>

Masses for gauginos of the group G_i arising in the D7$_i$-D7$_i$ sector are denoted M_i or M_{G_i}. Then,

$$\begin{aligned} M_1 &= M_{SU(4)} = \sqrt{3} C m_{3/2} \left[e^{-i\gamma_S} \sin\theta \sin^2\pi\delta + e^{-i\gamma_1} \eta_1 \cos\theta \cos^2\pi\delta \right] \ , \\ M_2 &= M_{SU(2)_L} = \sqrt{3} C m_{3/2} e^{-i\gamma_2}\, \eta_2 \cos\theta \ , \\ M_3 &= M_{SU(2)_R} = \sqrt{3} C m_{3/2} e^{-i\gamma_3}\, \eta_3 \cos\theta \ . \end{aligned} \tag{4.5}$$

<u>Scalar masses of SM fields</u>.

The Higgs multiplets appear in a twisted sector (unmagnetized) from D7$_2$-D7$_3$ intersections. One finds then the simple result

$$m_H^2 = m_{3/2}^2 + V_0 - \frac{3}{2} C^2 m_{3/2}^2 \left(\sin^2\theta + \eta_1^2 \cos^2\theta \right) \ . \tag{4.6}$$

On the other hand, the chiral quark/lepton fields appear on magnetized and twisted sectors D7$_1$-D7$_2$ and D7$_1$-D7$_3$. In particular, all three generations of left-handed quarks and leptons come from the D7$_1$-D7$_2$ sector and have soft masses

$$\begin{aligned} m_{L_\alpha}^2 &= m_{3/2}^2 + V_0 - \frac{3}{2} C^2 m_{3/2}^2 \left(\sin^2\theta + \eta_3^2 \cos^2\theta \right) + \frac{3}{16\pi^2} C^2 m_{3/2}^2 \sin^2 2\pi\delta B_1(\delta) |\Theta|^2 \\ &+ \frac{3}{8\pi} C^2 m_{3/2}^2 \sin 2\pi\delta \left[2 \left(\sin^2\theta - \eta_1^2 \cos^2\theta \right) - \cos 2\pi\delta |\Theta|^2 \right] \left[\log\frac{u_3}{u_2} + B_0(\delta) \right] \ , \end{aligned} \tag{4.7}$$

where we have defined

$$
\begin{aligned}
\Theta &= e^{-i\gamma_S}\sin\theta - e^{-i\gamma_1}\eta_1\cos\theta \ , \\
B_0(\delta) &= \psi_0(1-\delta) - \psi_0(\tfrac{1}{2}-\delta) \ , \\
B_1(\delta) &= \psi_1(1-\delta) - \psi_1(\tfrac{1}{2}-\delta) \ .
\end{aligned}
\tag{4.8}
$$

Here $\psi_0(z) = \Gamma'(z)/\Gamma(z)$ and $\psi_1(z) = \psi_0'(z)$.

The three generations of right-handed quarks and leptons come from the magnetized and twisted D7_1-D7_3 sector. The soft masses $m^2_{R_\alpha}$ have the same form as (4.7) except for the replacements $u_2 \leftrightarrow u_3$, $\eta_2 \leftrightarrow \eta_3$. Note that the limit $\delta \to 0$ just yields the scalar mass for unmagnetized twisted fields like the Higgs multiplet, as it should. This also agrees with the results obtained for unmagnetized branes in [26].

As a check on the results we can use

$$
\psi_0(\tfrac{1}{2}-\delta) = \psi_0(\tfrac{1}{2}+\delta) - \pi\tan\pi\delta \ , \quad ; \quad \psi_1(\tfrac{1}{2}-\delta) = \pi^2\sec^2\pi\delta - \psi_1(\tfrac{1}{2}+\delta) \ , \tag{4.9}
$$

to take the limit $\delta \to \frac{1}{2}$. This corresponds to infinite magnetic flux in the 2nd and 3rd torus. In this limit the magnetized D7_1-brane behaves as a D3-brane. Thus, taking $\delta \to \frac{1}{2}$ in (4.7) should give the mass squared parameter of a scalar in a D3-D7_2 type of sector. In this way we obtain

$$
m^2_{C^{37_2}} = m^2_{3/2} + V_0 - \frac{3}{2}C^2 m^2_{3/2}(1-\eta_2^2)\cos^2\theta \ , \tag{4.10}
$$

in accordance with the expected outcome [26].

<u>Trilinear terms of SM fields</u>.

The coupling $HL_\alpha R_\beta$ is of type $C^{7_17_2}C^{7_27_3}C^{7_37_1}$. The trilinear term turns out to be

$$
A_{HLR} = \frac{\sqrt{3}}{2}Cm_{3/2}\left[e^{-i\gamma_S}\sin\theta - \sum_i e^{-i\gamma_i}\eta_i\cos\theta - \frac{1}{\pi}B_0(\delta)\sin 2\pi\delta\,\Theta\right] \ . \tag{4.11}
$$

When $\delta \to 0$, the result agrees with that in [26]. One can also take the limit when $\delta \to \frac{1}{2}$ which should correspond to a coupling of type $C^{37_2}C^{7_27_3}C^{37_3}$. It indeed follows that

$$
A_{C^{37_2}C^{7_27_3}C^{37_3}} = \frac{\sqrt{3}}{2}Cm_{3/2}\left[(e^{-i\gamma_1}\eta_1 - e^{-i\gamma_2}\eta_2 - e^{-i\gamma_3}\eta_3)\cos\theta - e^{-i\gamma_S}\sin\theta\right] \ , \tag{4.12}
$$

also in agreement with [26].

The other relevant trilinear coupling involving SM fields is of the form $C_1^{7_j} C^{7_2 7_3} C^{7_2 7_3}$, $j = 2, 3$. For those couplings one gets trilinear terms $A = -M_j$. In our case the only such coupling is XHH. Then,

$$A_{XHH} = -\sqrt{3} C m_{3/2} e^{-i\gamma_2} \, \eta_2 \cos\theta = -M_{SU(2)} \ . \tag{4.13}$$

<u>Soft terms of non-chiral fields</u>.

Together with the chiral MSSM-like spectrum, there are three chiral multiplets in the adjoint of the gauge group coming from untwisted D7_j-D7_j sectors. We set $\Phi_{ij} = C_i^{7_j}$ and recall that Φ_{jj} parametrizes the position of each D7_j-brane in transverse space, whereas the Φ_{ij} correspond to Wilson lines on the two complex dimensions inside the D7_j-brane worldvolume. For the unmagnetized Φ_{i2} we find

$$\begin{aligned} m_{12}^2 &= m_{3/2}^2 + V_0 - 3C^2 m_{3/2}^2 \eta_3^2 \cos^2\theta \ , \\ m_{22}^2 &= m_{3/2}^2 + V_0 - 3C^2 m_{3/2}^2 \sin^2\theta \ , \\ m_{32}^2 &= m_{3/2}^2 + V_0 - 3C^2 m_{3/2}^2 \eta_1^2 \cos^2\theta \ . \end{aligned} \tag{4.14}$$

For the Φ_{i3} there are analogous results. For the magnetized Φ_{i1} the masses are instead

$$\begin{aligned} m_{11}^2 &= m_{3/2}^2 + V_0 - 3C^2 m_{3/2}^2 (\sin^2\theta + \eta_1^2 \cos^2\theta) + |M_1|^2 \ , \\ m_{21}^2 &= m_{3/2}^2 + V_0 - 3C^2 m_{3/2}^2 \eta_2^2 \cos^2\theta \ , \\ m_{31}^2 &= m_{3/2}^2 + V_0 - 3C^2 m_{3/2}^2 \eta_3^2 \cos^2\theta \ . \end{aligned} \tag{4.15}$$

In all cases there is a sum rule

$$m_{1j}^2 + m_{2j}^2 + m_{3j}^2 = 2V_0 + |M_j|^2 \ . \tag{4.16}$$

Besides scalar masses there are also trilinear terms associated to the superpotential couplings $\Phi_{1i}\Phi_{2i}\Phi_{3i}$. They are given by $A_i = -M_i$.

<u>General structure of soft terms</u>.

Clearly, the structure of soft terms strongly depends on which auxiliary fields, either F_{T_i} or F_S, dominate SUSY-breaking. As a general property one must emphasize that in all cases the results for scalar masses are flavor independent. The trilinear terms involving squarks and sleptons are also flavor diagonal. This is an interesting property which was not always present in heterotic models and is welcome in order to suppress too large flavor changing neutral currents (FCNC). Concerning gaugino masses, they are equal for the $SU(2)_L \times SU(2)_R$ sector of the theory but different for the $SU(3+1)$ gauginos which get an extra contribution proportional to F_S due to the presence of magnetic flux in the $D7_1$-brane.

One can easily check that in the diluted magnetic flux limit the results for soft terms agree with those found in [26]. Of particular interest are the extremes in which either the overall modulus T or the dilaton S auxiliary field dominate SUSY-breaking. T-dominance ($\cos\theta = 1$) appears for SUSY-breaking induced by ISD fluxes and we will discuss it in detail in the next section. Concerning dilaton dominance ($\sin\theta = 1$), eq. (4.7) shows that it is potentially dangerous since squarks and sleptons typically become tachyonic for small δ. However, we already mentioned that to stay within the string perturbative regime the value of δ cannot be too small, c.f. eq.(3.24). Thus, dilaton dominance could still be consistent if magnetic fluxes are substantial, in fact for $\delta \gtrsim 0.258$ when $u_2 = u_3$. One may argue that there are scalars whose masses become tachyonic in the limit $\sin\theta = 1$ and have no dependence on δ, since they are related to unmagnetized branes. This is the case of the Higgs and the non-chiral scalars Φ_{22} and Φ_{33} parameterizing the position of branes $D7_2$, $D7_3$. However, for both there could be extra contributions to their masses which would render them non-tachyonic. In the case of the Higgs multiplet we already mentioned that they get an additional SUSY mass term for $\langle X \rangle \neq 0$. Concerning the Φ_{22}, Φ_{33} scalars, they may also have SUSY μ-terms. For example, in flux induced SUSY-breaking, fluxes of type $(2,1)$ or $(1,2)$ could give rise to such terms (see [20]). All in all, whereas T-dominance always leads to a non-tachyonic structure of soft terms, in the case of dilaton dominance substantial magnetic fluxes and additional positive contribution for Higgsses and some adjoints are required to avoid tachyons.

4.1 Soft terms induced by fluxes

We now wish to discuss the situation in which the moduli superpotential is just given by W_f, c.f. (3.29). Then, fluxes are the only source of SUSY breaking.

Soft terms from T-dominance (ISD fluxes).

A concrete realization of T-dominance arises when the flux G_3 is generic ISD. In this case, $\langle F^S \rangle = \langle F^{U_i} \rangle = 0$ and only $\langle F^{T_i} \rangle \neq 0$. The cosmological constant vanishes automatically since W_f is independent of T_i and $\hat{K}$ is of no-scale form. Thus, $V_0 = 0$ and the other relevant vevs are

$$m_{3/2}^2 = \frac{|W_f|^2}{P} \quad ; \quad \bar{F}^{\bar{T}_i} = -\frac{t_i W_f}{P^{1/2}} = e^{i\gamma_T}\, t_i\, m_{3/2} \ , \tag{4.17}$$

where $P = s\prod_i t_i u_i$. Comparing with (4.3) and (4.4) shows that $C = 1$, $\cos\theta = 1$, $\eta_i = 1/\sqrt{3}$, $\gamma_i = \gamma_T$, $\forall i$. The soft terms follow substituting these values in the general expressions. The results are collected in Table 2 [2].

It is straightforward to expand the soft terms near $\delta = 0$. For example,

$$\begin{aligned} m_{L_\alpha}^2 &= m_{3/2}^2 \left(\frac{1}{2} - \frac{3}{4}\delta \log \frac{4u_3}{u_2} - \frac{\pi^2}{3}\delta^2 + \cdots \right) \ , \\ A_{HL_\alpha R_\beta} &= e^{-i\gamma_T}\, m_{3/2} \left(-\frac{3}{2} + 2\delta \log 2 + \frac{\pi^2}{3}\delta^2 + \cdots \right) \ . \end{aligned} \tag{4.18}$$

Using (4.9) we can also take the limit $\delta \to \frac{1}{2}$ in which $\mathrm{D7}_1 \to \mathrm{D}_3$. We find, $m_{L_\alpha}^2 \to 0$ and $A_{HL_\alpha R_\beta} \to -\frac{1}{2}e^{-i\gamma_T}\, m_{3/2}$, matching results of [26].

Note that the structure of soft terms in this subsection corresponds to SUSY-breaking induced by the auxiliary field of the overall Kähler modulus T, and the fact that it may be induced by fluxes plays no role in the obtained results. A few comments on the structure of soft terms in this simple case are in order.

- The structure of scalar soft terms is not universal, i.e. the scalar masses of left-handed sfermions, right-handed sfermions and Higgsses are different. However they are *flavor independent*. This is due to the origin of family replication in this class

[2] When the ISD flux has both (2,1) and (0,3) components there is an induced supersymmetric μ term and an extra soft bilinear parameter for the Φ_{ii} scalars [20].

$m^2_{L_\alpha}$	$\frac{1}{2} \ - \frac{1}{8\pi}\sin 2\pi\delta(2+\cos 2\pi\delta)[\log\frac{u_3}{u_2} + B_0(\delta)] + \frac{1}{16\pi^2}\sin^2 2\pi\delta B_1(\delta)$
$m^2_{R_\beta}$	$\frac{1}{2} \ - \frac{1}{8\pi}\sin 2\pi\delta(2+\cos 2\pi\delta)[\log\frac{u_2}{u_3} + B_0(\delta)] + \frac{1}{16\pi^2}\sin^2 2\pi\delta B_1(\delta)$
m^2_H	$\frac{1}{2}$
$M_{SU(3+1)}$	$e^{-i\gamma_T}\cos^2\pi\delta$
$M_{SU(2)_L}$	$e^{-i\gamma_T}$
$M_{SU(2)_R}$	$e^{-i\gamma_T}$
$A_{HL_\alpha R_\beta}$	$e^{-i\gamma_T}\ \left[-\frac{3}{2}\ + \frac{1}{2\pi}\sin 2\pi\delta B_0(\delta)\right]$
A_{XHH}	$-e^{-i\gamma_T}$
$m^2_{\Phi_{jj}}$	$\lvert M_j\rvert^2$
$m^2_{\Phi_{ij}}$	0

Table 2: Soft terms for T-dominant ISD fluxes. Results are given in $m_{3/2}$ units.

of models. The massless chiral fermions come in identical replicas with diagonal kinetic terms. Concerning gaugino masses, as we said the $SU(2)_L$ and $SU(2)_R$ gaugino masses are identical but that of $SU(3+1)$ is different.

- Note that if at the end of the day a non-vanishing FI-term (2.8) is present, an extra contribution to squark/slepton masses with opposite signs for left and right-handed fields will be added. This contribution will not be present for the Higgs fields which are neutral under the anomalous $U(1)$.

- In the formal limit $\delta \to 0$ corresponding to diluted magnetic fluxes one obtains particularly simple and universal results for soft terms :

$$\begin{aligned} m^2_{L_\alpha} &= m^2_{R_\alpha} = m^2_H = \frac{1}{2}m^2_{3/2} \qquad (4.19)\\ M_{SU(3+1)} &= M_{SU(2)_L} = M_{SU(2)_R} = e^{-i\gamma_T}\ m_{3/2}\\ A_{HL_\alpha R_\beta} &= -\frac{3}{2}m_{3/2}e^{-i\gamma_T}\\ A_{XHH} &= e^{-i\gamma_T}\ m_{3/2} \qquad (4.20)\end{aligned}$$

These results correspond to those advanced in section 6 of [21] for $\mu = \langle X \rangle$ and $\xi = 1/2$.

IASD fluxes and dilaton dominance.

Let us now analyze the mass parameters generated by the presence of IASD fluxes. Some words of caution should first be given. Care should be taken in comparing the results below for IASD fluxes to those for dilaton dominance ($\sin\theta = 1$) in the previous section. Indeed in the analysis of that section the values of V_0, $\sin\theta$, $m_{3/2}$, appear as independent parameters. Thus, in principle one can conceive a situation with $\sin\theta = 1$, $V_0 = 0$ and $m_{3/2} \neq 0$ with SUSY broken in Minkowski space. However, with IASD (3,0) fluxes one can show that $V_0 \neq 0$ and $m_{3/2} = 0$, so we have broken SUSY in de Sitter space. We will still provide the generated mass terms for completeness.

If we consider W_f as our only source of SUSY-breaking dilaton dominance appears when the flux G_3 is IASD of (3,0) type. In the absence of a W_{np} term this background is not in general a solution of the Type IIB equations of motion. However, it is a simple example of S-dominance because $\langle F^{T_i} \rangle = \langle F^{U_i} \rangle = 0$ but $\langle F^S \rangle \neq 0$. In this case $m_{3/2} = 0$ automatically and, as we said, there is a cosmological constant $V_0 \neq 0$. In terms of $Y_f = \int \bar{G}_3 \wedge \Omega$ one finds

$$V_0 = \frac{|Y_f|^2}{P} \quad ; \quad \bar{F}^{\bar{S}} = -\frac{sY_f}{P^{1/2}} = e^{i\gamma_S}\, s\sqrt{V_0}\ . \tag{4.21}$$

Hence, the soft terms can be obtained from the general expressions setting $\sin\theta = 1$ and $Cm_{3/2} \to \sqrt{V_0/3}$. Results are displayed in Table 3. Note that the mass parameters are in general not tachyonic. This is not in contradiction with our results discussed at the end of section 3, since there we assumed arbitrary V_0 and $m_{3/2} \neq 0$, whereas IASD $(3,0)$ fluxes lead to $V_0 \neq 0$ and $m_{3/2} = 0$.

When ISD and IASD fluxes are turned on simultaneously the auxiliary fields, $m_{3/2}$ and V_0 are just given by (4.17) and (4.21). Then, $\eta_i = 1/\sqrt{3}$, $\gamma_i = \gamma_T$, $\forall i$, $3\tan^2\theta = V_0/3m_{3/2}$, and $C = \sec\theta$. The soft terms can be found substituting these values in the general expressions. In most cases it suffices to add the entries in Tables 2 and 3.

Let us end this section with some comments concerning the mass scales in this model. Note that as it stands, in a *toroidal* model like this, the string scale M_s should be of order

$m^2_{L_\alpha}$	$\frac{1}{2} \ + \frac{1}{8\pi}\sin 2\pi\delta(2-\cos 2\pi\delta)[\log\frac{u_3}{u_2} + B_0(\delta)] + \frac{1}{16\pi^2}\sin^2 2\pi\delta B_1(\delta)$
$m^2_{R_\beta}$	$\frac{1}{2} \ + \frac{1}{8\pi}\sin 2\pi\delta(2-\cos 2\pi\delta)[\log\frac{u_2}{u_3} + B_0(\delta)] + \frac{1}{16\pi^2}\sin^2 2\pi\delta B_1(\delta)$
m^2_H	$\frac{1}{2}$
$M_{SU(3+1)}$	$e^{-i\gamma_S}\sin^2\pi\delta$
$M_{SU(2)_L}$	0
$M_{SU(2)_R}$	0
$A_{HL_\alpha R_\beta}$	$e^{-i\gamma_S} \ \left[\frac{1}{2} \ - \frac{1}{2\pi}\sin 2\pi\delta B_0(\delta)\right]$
A_{XHH}	0
$m^2_{\Phi_{jj}}$	$\lvert M_j\rvert^2$
$m^2_{\Phi_{ij}}$	1

Table 3: Soft terms for S-dominant IASD fluxes. Results are given in V_0 units.

(or slightly smaller) than the Planck scale. Indeed, both scales are related by eq.(3.5). Although one may think that one can make $M_{Pl} >> M_s$ by taking t_i very large, that would make the SM gauge couplings unacceptably small, as shown by eqs.(3.22). On the other hand, the size of SUSY-breaking soft terms depends on the value of the gravitino mass in these theories. In general, if the source of SUSY-breaking is not specified, as in section 3, one can assume that $m_{3/2}$ may be small, i.e. of order the electroweak scale, as in the canonical approach to gravity mediated SUSY-breaking models. This was our general philosophy in the first part of section 4. On the other hand, if one insists that the microscopic source of SUSY-breaking is some ISD flux in a *toroidal* setting , then the gravitino mass is given by eq.(4.17). In that case, since the t_i fields cannot be too large, the gravitino mass is of order the string scale (which is only slightly smaller than M_{Pl}) and hence too large to lead to a solution of the hierarchy problem. This is the fact already mentioned in the introduction. However, as we said, this is a particular property of toroidal settings in which fluxes are distributed uniformly in extra dimensions. One can conceive an embedding of the MSSM-like brane setting in [1] into a CY/F-theory

compactification in which the distribution of fluxes in extra dimensions is not constant and hierarchically small soft terms may appear.

5 Conclusions

In this paper we have computed the SUSY-breaking soft terms for the MSSM-like model introduced in [1] under the assumption of generic vevs for the auxiliary fields F_{T_i} and F_S. We provide the soft terms as explicit functions of the gravitino mass, goldstino angle and a parameter δ that characterizes the magnetic flux in one of the brane stacks. We find that the case of isotropic T-dominance is particularly interesting since it always leads to simple results with no tachyons. For dilaton dominance there is the risk of getting some tachyonic masses for SM fields unless magnetic fluxes are large and additional sources for masses of non-chiral fields are present. The case of isotropic T-dominance appears in particular when SUSY-breaking is triggered by ISD antisymmetric Type IIB fluxes. We argue that although in a toroidal setting the soft terms induced by fluxes are typically too large, they may be hierarchically small in more general CY/F-theory embeddings of this MSSM-like brane configuration.

The results for soft terms in T-dominance are summarized in Table 2, and take an even simpler form (4.20) in the dilute flux limit $\delta \to 0$. They are flavor universal and depend only on the values of $m_{3/2}$, δ and a complex phase γ_T.

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